

Electroweak Symmetry Breaking Beyond the Standard Model ¹

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Abstract

In this talk, I shall address two key issues related to electroweak symmetry breaking. First, how fine-tuned different models are that trigger this phenomenon? Second, even if a light Higgs boson exists, does it have to be necessarily elementary? After a brief introduction, I shall first review the fine-tuning aspects of the MSSM, NMSSM, generalized NMSSM and GMSB scenarios. I shall then compare and contrast the little Higgs, composite Higgs and the Higgsless models. Finally, I shall summarize by giving a broad overview on where we stand at the end of 2011.

I Introduction

The timing of the last ‘Lepton-Photon Conference’ (August 2011) was very special! Every day the LHC was delivering more data than it did during the entire 2010. The time for ‘speculation’ was soon coming to an end! Our imagination about the possible dynamics behind electroweak symmetry breaking (EWSB), disciplined by the constraints from electroweak precision tests (EWPT), has fueled different directions of theoretical studies and experimental searches over the last so many years. Finally, the LHC has roared into life, and this is our last chance of putting money on our favorite models. It is in this backdrop that I have prepared a write-up of my talk, being aware that even during the last few months since Lepton-Photon the excluded territory for different Beyond the Standard Model (BSM) alternatives has further grown in size.

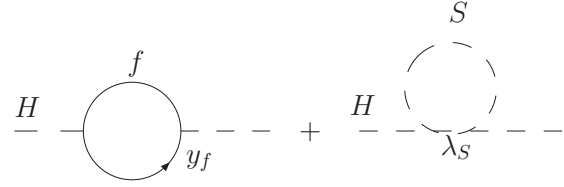


Figure 1: Cancellation of quadratic divergence to scalar mass-square between fermion and boson loops.

Now, to the point. We know that the SM Higgs mechanism is only an effective description of EWSB. Can LHC shed enough light on the dynamics behind this mechanism? Some of the questions that drive our speculation are listed below [1]:

- (i) Why is the weak scale so much separated from the Planck scale?
- (ii) What is the symmetry that controls particle physics at the TeV scale? In other words, now that the gauge symmetry is established with a significant precision, what is the next relevant symmetry that awaits us?
- (iii) The SM is plagued by the hierarchy problem. It originates from the requirement of *ad hoc* cancellation between fermionic and bosonic loops contributing to the Higgs mass – see Fig. 1. An unnatural tuning ($1 \div 10^{26}$) between the bare Higgs mass-square $m_{h_0}^2$ and the correction term Δm_h^2 is necessary to keep the renormalized mass ($m_h^2 = m_{h_0}^2 + \Delta m_h^2$) at around 100 GeV. Nevertheless, one must do this tuning order-by-order in perturbation theory to prevent the Higgs mass

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from shooting up to the highest scale of the theory. This constitutes the hierarchy problem. Quite a few remedies have been advocated so far. But, which solution (if any, at all!) of the hierarchy problem is correct?

(iv) Is the naturalness consideration a good guiding principle or a powerful discriminator between models? Is its study a step in the right direction [2]?

(v) Is Higgs *elementary* or *composite* [3, 4]? Can it be settled at the LHC?

(vi) What if the Higgs is not there at all?

II Supersymmetry

II.1 Basic aspects

Supersymmetry is the most well-studied BSM model that offers a natural explanation of the weak scale [5]. It is a new space-time symmetry interchanging bosons and fermions, relating states of different spins. The Poincare group is extended by adding two anti-commuting generators Q and \bar{Q} to the existing p (linear momentum), J (angular momentum) and K (boost), such that $\{Q, \bar{Q}\} \sim p$. Since the new symmetry generators are spinors, not scalars, supersymmetry is not an internal symmetry, and the super-partners differ from their SM partners in spin. Some attractive features of supersymmetry relevant in the present context are as follows:

(i) *Supersymmetry solves the gauge hierarchy problem*: The quantum corrections to the Higgs mass from a bosonic loop and a fermionic loop exactly cancel if the couplings are identical and the boson is mass degenerate with the fermion. For every fermion (boson) of the SM, supersymmetry provides a mass degenerate boson (fermion). In real life, however, supersymmetry is badly broken. But if the breaking occurs in masses and not in dimensionless couplings, the quadratic divergence still cancels. The residual divergence is only logarithmically sensitive to the supersymmetry breaking scale.

(ii) *Supersymmetry leads to gauge coupling unification*: This is a bonus! Supersymmetry was not invented to achieve this. When the SM gauge couplings are extrapolated to high scale, with LEP measurements as input, they do not meet at a single point. Supersymmetry makes them do at a scale $M_{\text{GUT}} \sim 2 \times 10^{16}$ GeV, with TeV scale super-particles.

(iii) *Supersymmetry triggers EWSB*: Starting from a positive value in the ultraviolet, the up-type Higgs mass-square $m_{H_u}^2$ turns negative in the infrared triggering EWSB. In the SM the negative sign in front of the scalar mass-square in the potential is put in by hand to ensure EWSB. In supersymmetry the sign flip occurs in a dynamical way.

II.2 Naturalness criterion

Naturalness is an aesthetic criterion. It comes from the realization that if large cancellation among *unrelated* quantities is required to achieve a small physical quantity, the situation is unnatural and reflects a sign of *weak health* of the theory. A theory is less ‘natural’ if it is more ‘fine-tuned’. In the context of minimal supersymmetry with two Higgs doublet (i.e. MSSM), the scalar potential minimization yields

$$\frac{1}{2}M_Z^2 = \frac{m_{H_d}^2 - m_{H_u}^2 \tan^2 \beta}{\tan^2 \beta - 1} - \mu^2, \quad (1)$$

with $m_{H_u}^2 = m_{H_d}^2 - \Delta m^2$, where Δm^2 is the correction due to RG running from the GUT scale to the weak scale. The large top Yukawa coupling has a significant numerical influence on RG running. A proper EWSB occurs when $m_{H_u}^2$ turns negative due to the effect of running and the correct value of M_Z is reproduced. This refers to a cancellation between

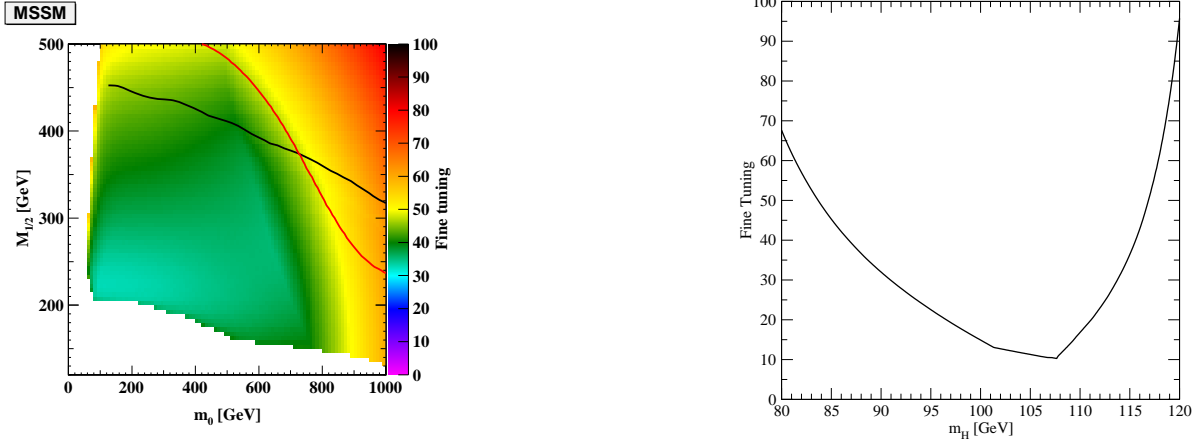


Figure 2: (a) Left: Fine-tuning in cMSSM. (b) Right: Fine-tuning as a function of m_h . For details see [8].

supersymmetry breaking soft masses and supersymmetry preserving μ parameter. How much cancellation between these completely uncorrelated quantities is aesthetically pleasant? Barbieri and Giudice introduced a quantitative measure of fine-tuning [6]

$$\Delta_i \equiv \left| \frac{\partial M_Z^2 / M_Z^2}{\partial a_i / a_i} \right|, \quad (2)$$

where a_i are high scale input parameters. An upper limit on Δ can be translated to an upper limit on super-particle masses.

II.3 Naturalness of cMSSM

In cMSSM, the constrained version of the MSSM (with 4 parameters and 1 sign), Eq. (1) boils down to [7]

$$M_Z^2 \approx -2|\mu^2| + 0.2 m_0^2 + 0.7 (2.6 M_{1/2})^2, \quad (3)$$

where m_0 and $M_{1/2}$ are the common scalar and gaugino masses, respectively, and the gluino mass is given by $m_{\tilde{g}} \simeq 2.6 M_{1/2}$. Two observations are noteworthy:

(i) In the absence of any cancellation, the natural expectation would be $M_Z \sim \mu \sim m_0 \sim M_{1/2}$. But this possibility has been explored and ruled out by LEP-2 and Tevatron.

(ii) By now, the CMS and ATLAS collaborations have pushed the gluino mass limit to close to a TeV. This implies a tuning of order 1% from Eq. (3). The LHC is thus probing sparticle masses which are about a loop factor above M_Z .

There is another way to show that the fine-tuning in MSSM is $\sim 1\%$ level. The radiatively corrected mass of the lightest CP-even Higgs boson is given by

$$m_h^2 \simeq m_{h0}^2 (\leq M_Z^2) + \frac{3m_t^4}{2\pi^2 v^2} \ln \left(\frac{m_{\tilde{t}}^2}{m_t^2} \right). \quad (4)$$

Since $m_h > 114$ GeV from LEP-2, $m_{\tilde{t}}$ should be around 1 TeV or heavier, thus implying a fine-tuning to the tune of a per-cent. This constitutes the ‘little hierarchy’ problem of supersymmetry.

A quantitative analysis of fine-tuning has recently been carried out in [8] (see also [9] where some of the technical aspects for measuring the tuning are a little different) in the context of the cMSSM. Fig. 2a corresponds to $\tan \beta =$

3 and $A_0 = 0$. The different parts of the white region is ruled out for different reasons (non-occurrence of EWSB, experimental exclusion of the slepton/neutralino/chargino mass limits, Higgs mass lower limit, stau becoming the LSP). The experimental bounds from ATLAS (black) and CMS (red) have been drawn for a guide to the eye using $1/fb$ data. The fine-tuning is at best $\sim 2\%$ which corresponds to $\Delta \sim 50$. We now look at Fig. 2b where fine-tuning has been plotted against m_h . The LEP-2 lower limit has not been imposed here. It is interesting to see that the tuning is minimum around $m_h = 108$ GeV. If m_h is lower than that, the fine-tuning becomes larger as sparticle masses are constrained by their experimental lower limits. On the other hand, if m_h is higher than this value, then due to its $\ln(m_{\tilde{t}})$ -dependence there is an exponential growth of fine-tuning.

It is interesting to note that for values of $m_0 \leq 700$ GeV and $M_{1/2} \leq 350$ GeV, the amount of fine-tuning is decided by the LEP-2 limit on the Higgs mass. On the other hand, for larger m_0 and $M_{1/2}$, the origin of fine-tuning can be traced to the adjustment between μ^2 and scalar soft mass-squares that yields the correct M_Z .

II.4 Naturalness of NMSSM

First we consider the NMSSM scenario which has an additional gauge singlet superfield S compared to MSSM [10]. The NMSSM superpotential has two important additional pieces,

$$W_{\text{NMSSM}} = W_{\text{Yukawa}} + \lambda S H_u H_d + \frac{1}{3} \kappa S^3. \quad (5)$$

The vev s of the scalar component of S yields an effective $\mu_{\text{eff}} = \lambda s$. In fact, this was the main motivation behind adding the singlet. The NMSSM models are less fine-tuned than MSSM for three reasons [8]:

(i) The $S H_u H_d$ term in Eq. (5) generates a quartic interaction in the scalar potential, increasing tree level m_{h_0} [11],

$$m_{h_0}^2 \approx M_Z^2 \cos^2 2\beta + \lambda^2 v^2 \sin^2 2\beta. \quad (6)$$

The additional tree level contribution allows us to consider a lighter stop in the loop to generate the same Higgs mass as in the MSSM. Fine-tuning is therefore reduced.

(ii) The physical Higgs boson can have a large singlet admixture, and therefore, a reduced gauge coupling which helps it evade the LEP-2 limit. This again implies that we can employ a lighter stop in the loop, thus reducing fine-tuning.

(iii) The possibility of Higgs decaying into two lighter pseudo-scalars also helps to evade the LEP-2 limit.

The minimal fine-tuning in NMSSM has been plotted in Fig. 3a. For smaller values of m_0 and $M_{1/2}$, i.e. in the region where the LEP-2 limit on m_h is the relevant constraint, fine-tuning is considerably less than in cMSSM. Δ can be as small as 14 in this region (as against 33 for cMSSM). However, for larger values of m_0 and $M_{1/2}$ the origin of fine-tuning lies in the smallness of weak scale compared to the soft masses, and in this region it is hard to reduce fine-tuning. Overall, NMSSM is less fine-tuned than cMSSM, or for that matter in MSSM with universal boundary conditions.

II.5 Naturalness of Generalized NMSSM (G-NMSSM)

G-NMSSM has an underlying Z_4 or Z_8 discrete symmetry [12]. Its superpotential reads

$$W_{\text{G-NMSSM}} = W_{\text{Yukawa}} + (\mu + \lambda S) H_u H_d + \frac{1}{2} \mu_S S^2 + \frac{1}{3} \kappa S^3, \text{ where } \mu \sim \mu_S \sim \mathcal{O}(m_{3/2}). \quad (7)$$

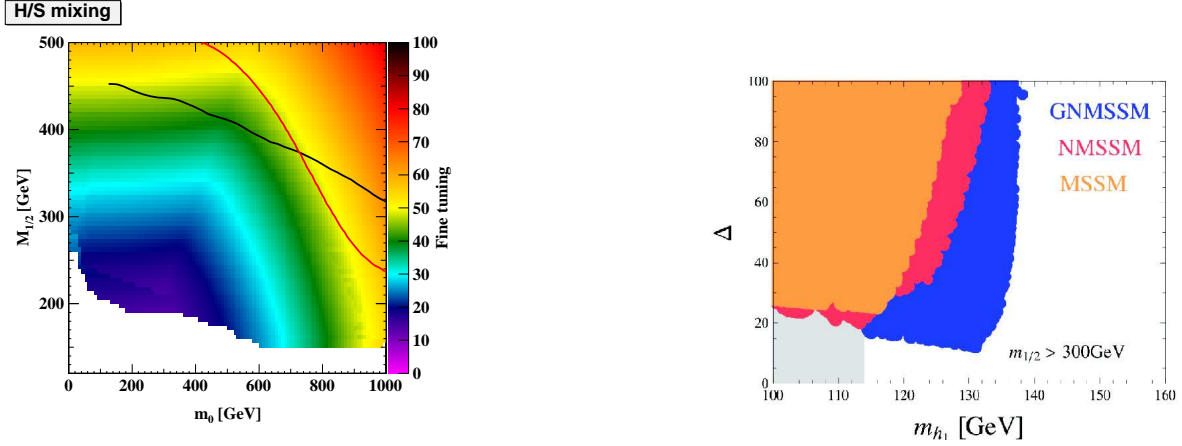


Figure 3: (a) Left: Fine-tuning in NMSSM [8]. (b) Right: Fine-tuning in G-NMSSM [12].

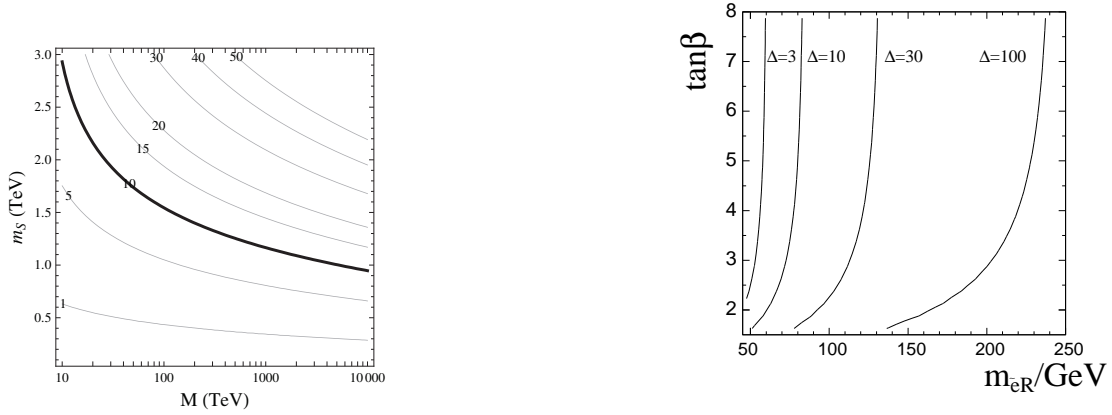


Figure 4: (a) Left: Fine-tuning in λ SUSY [14]. (b) Right: Fine-tuning in GMSB [15].

It has two distinct advantages beyond NMSSM. First, it has a discrete R symmetry, contrary to a discrete – but non- R – symmetry in NMSSM, which helps to remove the domain wall problems present in NMSSM [13]. The R symmetry is broken at a very high scale making the domain walls decay well before nucleosynthesis. And secondly, fine-tuning in G-NMSSM is considerably less than in NMSSM. The main reason behind this is the additional stabilizing terms in the potential. To appreciate this, take a large μ_S limit and integrate out the S superfield at the supersymmetric level. This gives a term $\lambda^2 (H_u H_d)^2 / \mu_S$ in the superpotential, which reduces fine-tuning. For a fixed value of Δ one gets a heavier Higgs and, interestingly enough, the fine-tuning is minimum for $m_h \sim 130$ GeV. Fig. 3b shows us how fine-tuning improves from MSSM to NMSSM and from NMSSM to G-NMSSM.

II.6 Naturalness of λ SUSY

Consider the NMSSM and assume that the trilinear coupling λ is rather large [14], at least ~ 1 at the weak scale. The sole purpose here is to reduce fine-tuning by increasing the singlet-induced tree level contribution to the Higgs mass, so that the dominant term is $m_{h0}^2 \sim \lambda^2 v^2 \sin^2 2\beta$. For example, the values $m_h^{\max} \simeq 2(3)M_Z$ for $\lambda(L) = \sqrt{4\pi}$ correspond to $\Lambda = 10^4$ TeV (100 TeV) – see Fig. 4a. The flip side is that by having such a low cutoff, the prized possession of supersymmetry, namely, gauge coupling unification, is sacrificed to buy naturalness!

II.7 Naturalness of GMSB models

In gauge mediated supersymmetry breaking (GMSB) models, a natural determination of M_Z in terms of the model parameters yields a rather upper limit on the mass of the right selectron – see Fig. 4b. The universal boundary conditions of scalar masses in cMSSM do not permit the lightest scalar to be much lighter than $m_{H_u}^2$. But in gauge mediated models, the proportionality of scalar masses to different gauge couplings (square) at the messenger scale M creates quite a bit of splitting among the different scalar masses at the weak scale, which in turn leads to *more* fine-tuning than in cMSSM [15].

III Little Higgs

III.1 Basic aspects

Little Higgs models were introduced as a solution to the little hierarchy problem (for a review, see [16, 17]). The Higgs is considered to be a pseudo-Goldstone boson associated with some global symmetry breaking. A Goldstone boson ϕ has a shift symmetry $\phi \rightarrow \phi + c$, where c is a constant, and as long as this symmetry is maintained a Goldstone remains massless at all order. But if there is an interaction which couples ϕ not as $\partial_\mu \phi$ the shift symmetry is explicitly broken and the Goldstone becomes massive. This way we get a pseudo-Goldstone boson. Recall that pion is a Goldstone which results from the spontaneous breaking of chiral symmetry group $SU(2)_L \times SU(2)_R$ to the isospin group $SU(2)_I$. Since quark masses and electromagnetic interaction explicitly break the chiral symmetry, pions are in fact pseudo-Goldstone bosons. Electromagnetism attributes a mass to π^+ of order $m_{\pi^+}^2 \sim (e^2/16\pi^2)\Lambda_{\text{QCD}}^2$. If we think of Higgs mass generation in the same way, using gauge or Yukawa interaction as a source for explicit breaking of the chiral symmetry, we can have $m_h^2 \sim (g^2/16\pi^2)\Lambda_{\text{NP}}^2$. This picture is not phenomenologically acceptable, since $m_h \sim 100$ GeV implies $\Lambda_{\text{NP}} \sim 1$ TeV, but such a low cutoff is strongly disfavored by EWPT. If, on the other hand, we can somehow arrange that the leading term in the Higgs mass is

$$m_h^2 \sim \frac{g_1^2 g_2^2}{(16\pi^2)^2} \Lambda_{\text{NP}}^2, \quad (8)$$

then for a 100 GeV Higgs mass, we get $\Lambda_{\text{NP}} \sim 10$ TeV. The cutoff is thus postponed from 1 to 10 TeV thanks to the extra suppression factor of $16\pi^2$, without having to apparently pay any price for fine-tuning. The idea of ‘little Higgs’ is all about achieving this extra $16\pi^2$ factor in the denominator of Eq. (8), and this is where it differs from a pion. Note that both g_1 and g_2 should be simultaneously non-vanishing in order to generate the Higgs mass. If any of these couplings vanishes then the global symmetry is partially restored and the Higgs remains a Goldstone boson. This is the concept of ‘collective symmetry breaking’.

The basic features of the little Higgs trick are depicted in Fig. 5a. The global group G spontaneously breaks to H at a scale $f(> v)$. A part of G , labeled F , is weakly gauged and the overlap region between F and H is the unbroken SM group I . The Higgs, which is a doublet of the gauged $SU(2)$ of the SM, is a part of the Goldstone multiplet that parametrizes the coset space G/H . The generators corresponding to Higgs do not commute with the heavy gauge boson generators. Gauge (also, Yukawa) interactions induce mass to the Higgs boson at one-loop level. Since the gauge group is expanded, we have additional gauge bosons and fermions. The quadratic divergence to the Higgs mass at one-loop level arising from a Z boson loop cancels against a similar contribution from a heavy Z_H loop, and the same thing happens between a t loop and a heavy T loop – see Fig. 5b. This is an example of ‘same statistics cancellation’.

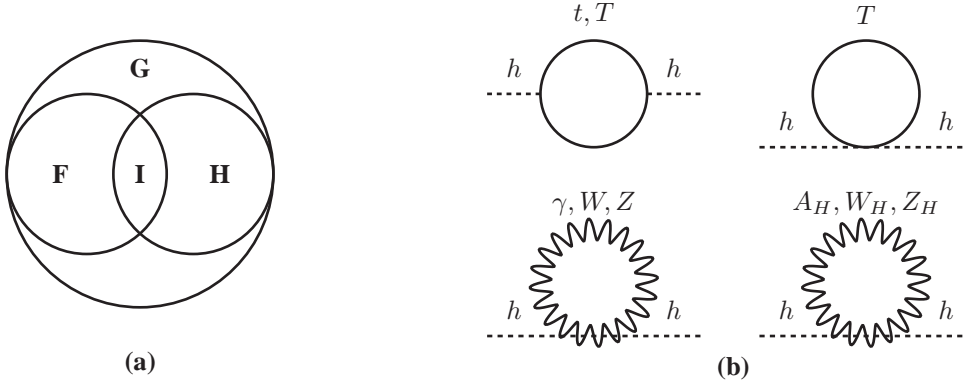


Figure 5: (a) Left: Little Higgs cartoon. (b) Right: Feynman diagrams among which same statistics cancellation takes place. T is a new heavy quark, and A_H, W_H, Z_H are new heavy gauge bosons.

III.2 Two crucial features

(i) The same statistics cancellation enables us to express $m_h^2 \sim f^2/16\pi^2 \ln(\mathbb{L}^2/f^2)$. But the quadratic cutoff sensitivity comes back parametrically at two-loop order. The order parameter is f is not protected from quadratic cutoff sensitivity, just like the electroweak vev v is not [18]. As a result,

$$f^2 \rightarrow F^2 = f^2 + \frac{\Lambda^2}{16\pi^2} \quad (\text{where } \mathbb{L} \sim 4\pi f), \quad (9)$$

Then, what did we gain compared to the SM? For little Higgs models

$$m_h^2 (\text{LH}) \sim \left(\frac{F^2}{16\pi^2} \right) \ln \left(\frac{\mathbb{L}^2}{F^2} \right). \quad (10)$$

This implies that $\Delta m_h^2 (\text{LH}) \sim \mathbb{L}^2/(16\pi^2)^2$, which should be compared with $m_h^2 (\text{SM}) \sim \mathbb{L}^2/(16\pi^2)$. For little Higgs, we thus have an extra suppression factor of $16\pi^2$, which indicates the parametric two-loop sensitivity of \mathbb{L}^2 . If we want $m_h \sim (f/4\pi) \sim 100$ GeV, one should have $f \sim F \sim 1$ TeV, and $\Lambda \sim 10$ TeV.

(ii) A clever construction of a little Higgs model should yield the following electroweak potential:

$$V = -\frac{(g_{\text{SM}})^4}{16\pi^2} f^2 \ln \left(\frac{\mathbb{L}^2}{f^2} \right) (H^\dagger H) + g_{\text{SM}}^2 (H^\dagger H)^2, \quad (11)$$

i.e., the bilinear term should have a one-loop suppression but, crucially, the quartic interaction should be *un-suppressed*, where g_{SM} is a gauge or Yukawa coupling. If both quadratic and quartic terms are suppressed, one cannot simultaneously obtain the correct W boson mass and an acceptable Higgs mass.

III.3 EWPT vs Naturalness

Contributions of new physics to two dimension-6 operators $\mathcal{O}_T \propto |H^\dagger D_\mu H|^2$ and $\mathcal{O}_S \propto H^\dagger \sigma^a H W_{\mu\nu}^a B_{\mu\nu}$ should be small enough to keep EWPT (T and S parameters, respectively) under control. A large class of little Higgs models gives a large contribution to T . Consequently, the constraint is quite strong: $f > (2 - 5)$ TeV [19]. A large

f means that to obtain the Higgs mass in the 100 GeV range one must fine-tune the parameters. The constraints arise primarily from the tree level mixing of the SM particles with the new particles. In the littlest Higgs model ($G = \text{SU}(5)$, $H = \text{SO}(5)$), the T parameter receives a large contribution from the custodial symmetry breaking operator $H^T \Phi H$, which mixes the doublet scalar H with the triplet scalar Φ . To avoid this mixing, the authors of [20] introduced T -parity (similar to R -parity in supersymmetry) under which all (*but one*) new particles are odd and the SM particles are even. Under this symmetry $H \rightarrow H$, but $\Phi \rightarrow -\Phi$, so $H^T \Phi H$ coupling is absent. As a result, f as low as 500 GeV can be accommodated [21]. Interestingly, there exists one new, yet T -even, state in this scenario, the so-called ‘top partner’, which cancels the standard top induced quadratic divergence to the Higgs mass.

Remember that we set out to solve the little hierarchy problem and, apparently, we settled that by acquiring an extra suppression factor of $16\pi^2$. But could we actually reduce the fine-tuning in realistic little Higgs models? Very importantly, a sizable tuning among various contributions to the Higgs quartic coupling is necessary to keep the Higgs mass small. Fine-tuning is relatively small when the Higgs mass is rather high, but this option is at odds with the requirement of EWPT. This underlines the tension between naturalness and EWPT. In fact, fine-tuning is $\leq 1\%$ in the phenomenologically acceptable region of the parameter space, and the general conclusion is that little Higgs models are less natural than MSSM [22] – see Fig. 6.

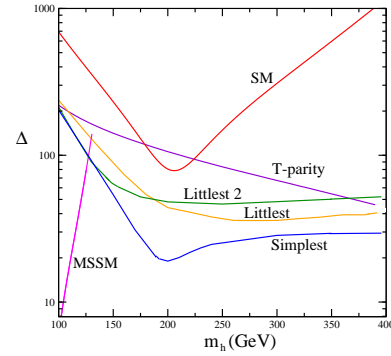


Figure 6: Fine-tuning in different little Higgs models (adapted from [22]).

III.4 Collider signals of little Higgs models

New gauge bosons: In the littlest Higgs model, about 30000 Z_H can be produced annually at the LHC with 100fb^{-1} luminosity. They would decay into the SM fermions ($Z_H \rightarrow f\bar{f}$), or into the SM gauge bosons ($Z_H \rightarrow W^+W^-$, $W_H \rightarrow WZ$, or into the Higgs and SM gauge boson ($Z_H \rightarrow Zh$). The branching ratios would follow a definite pattern, which would serve as the ‘smoking gun signals’ [23, 24].

New fermions: Colored vector-like T quark appears in almost all little Higgs models. It may be produced singly by $bW \rightarrow T$ at the LHC. Typically, $\Gamma(T \rightarrow th) \approx \Gamma(T \rightarrow tZ) \approx \frac{1}{2}\Gamma(T \rightarrow bW)$. These branching ratio relations would constitute a characteristic signature for T quark discovery [23, 25].

New scalars: The presence of a doubly charged scalar ϕ^{++} , as a component of a complex triplet scalar, is a hallmark signature of a large class of little Higgs models. Its decay into like-sign dileptons ($\phi^{++} \rightarrow \ell^+\ell^+$) would lead to an unmistakable signal with a separable SM background [23].

IV Composite Higgs

IV.1 Basic ideas

The composite Higgs models emerged as an improved realization of the little Higgs scenarios both in terms of UV completion and the naturalness consideration. In the composite picture the Higgs is some kind of a composite bound state emerging from a strongly interacting conformal sector [26] (for a review, see [3, 27] and references therein). It is a pseudo-Goldstone boson which results when a global group G of a strongly coupled sector breaks to H at a scale $f(> v)$. The coset G/H contains the Higgs. We know that AdS/CFT correspondence allows us to relate a strongly coupled 4d

theory to a weakly coupled 5d AdS theory. Using this correspondence, while on the CFT side the Higgs can be viewed as a pseudo-Goldstone of some strongly coupled dynamics, on the AdS side, in what is called the *Gauge-Higgs unification* scenario [28], the same Higgs can be interpreted as the 5th component of a gauge field (A_5) propagating in the warped extra dimension,

It is also called a *holographic Higgs* [29]. The holographic 5d to 4d translation involves the presence of two sectors - weak and strong. The weakly interacting sector containing elementary objects (the SM gauge bosons and some fermions) is located at the $y = 0$ (Planck) brane, and the strongly interacting CFT sector at the $y = L$ (TeV^{-1}) brane. The latter sector contains the TeV bound states at a scale $\sim 1/L$, and the Higgs is one such bound state. But to have a little hierarchy between m_h and $1/L$, we require the Higgs to be a Goldstone resulting from some $G \rightarrow H$ breaking in the CFT sector. More precisely, the Higgs is a pseudo-Goldstone as the couplings of the SM gauge and matter fields with the CFT sector *explicitly* break G .

A very satisfactory feature of composite Higgs is that a non-linearly realized global symmetry of the CFT sector protects its mass and guarantees the absence of quadratic divergence at all order. The finiteness of the Higgs mass can be understood as follows: the Higgs is at the TeV brane, the scalar that breaks the bulk gauge symmetry lives at the Planck brane. The Higgs mass is generated by radiative corrections with loops involving bulk KK gauge fields which propagate from one brane to another. This mediation mechanism involves a transmission of information from the Planck to TeV brane, which makes it a *non-local* effect, and hence the potential (and, therefore, the Higgs mass) so generated is calculable and finite. This is a big advantage over the conventional little Higgs construction which suffers from quadratic cutoff sensitivity at two loop level.

Also, the global symmetry that protects the Higgs mass is a symmetry of the strong CFT sector. Therefore, one expects to see a set of new electroweak resonances which should appear as *complete* multiplets of the global group. For example, in the $\text{SO}(5)/\text{SO}(4)$ model, additional fermionic states besides the SM fermions are required to fill the spinorial representation $\underline{4}$ of $\text{SO}(5)$. The spectrum of new particles can therefore reveal the nature of the global symmetry, and is certainly richer than that of the conventional little Higgs models.

IV.2 Collider tests

(i) A generic prediction of composite Higgs is that its gauge and the Yukawa couplings are reduced from their SM values [30]. It can be parametrized as ($\xi \equiv \frac{v^2}{f^2}$)

$$g_{hff} = g_{hff}^{\text{SM}} (1 - C_f \xi), \quad g_{hVV} = g_{hVV}^{\text{SM}} (1 - C_V \xi).$$

where $\xi \sim 0.2$ is small enough to keep the contribution of the new resonances to the oblique parameters under control. Here C_f and C_V are numbers which depend on the choices of the groups G and H . The question is, however, whether the Higgs production cross section times its branching ratios in different channels can be measured with an accuracy of, say, (10-20)% or better? We would perhaps need to go to super-LHC or better to ILC to confirm or rule out compositeness in a definitive way.

(ii) Since the gauge coupling of the Higgs is smaller than g , there will be incomplete cancellation of divergence in the gauge boson scattering amplitude. As a result,

$$A(V_L V_L \rightarrow V_L V_L) \sim s/f^2.$$

Therefore, one hopes to see excess events in $V_L V_L \rightarrow V_L V_L$ channels. Again, this discussion is not perhaps experimentally relevant before we reach 14 TeV, maybe not before the super-LHC stage!

5D Higgsless Model in Warped Space

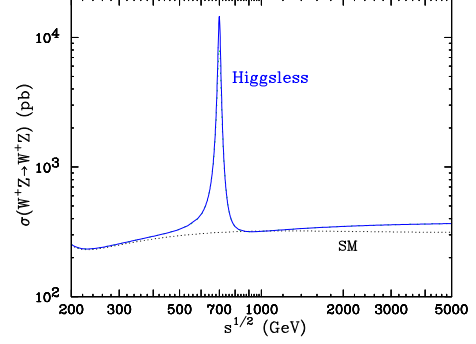
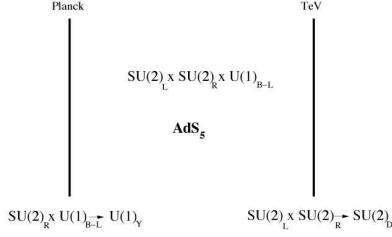


Figure 7: (a) Left: 5d Higgsless model [32], (b) Right: LHC signature in Higgsless model (Adapted from [35]).

(iii) The composite Higgs models usually contain heavy colored fermions of exotic charge, e.g. electric charge $5/3$, although this is a model dependent statement. Their production and decay may proceed as follows:

$$q\bar{q}, gg \rightarrow q_{5/3}^* \bar{q}_{5/3}^* \rightarrow W^+ t W^- \bar{t} \rightarrow W^+ W^+ b W^- W^- \bar{b}.$$

The decay products contain highly energetic same sign leptons, plus 6 jets two of which two are tagged b jets. Detecting those bound states would of course constitute the best test for compositeness [31].

V Higgsless scenarios

The idea here is to trigger electroweak symmetry breaking without actually having a physical Higgs. This is intrinsically an extra dimensional scenario. The basic construction goes as follows: the extra dimension is compactified on a circle of radius R with an orbifolding (S^1/Z_2). There are two fixed points: $y = 0, \pi R$. Electroweak breaking is achieved by imposing different boundary conditions (BC) on gauge fields at $y = 0, \pi R$. The BCs have to be carefully chosen such that the rank of a gauge group is lowered. The details can be found in [32]. The extra dimension can be flat or warped. It is difficult to control the T parameter in flat space, but in warped space one can construct a scenario which satisfies all EWPT constraints. Appropriate BCs are chosen to ensure the following gauge symmetry in the bulk and in the two branes (see Fig. 7a): Bulk: $SU(2)_L \times SU(2)_R \times U(1)_{B-L}$, $y = 0$ brane: $SU(2)_L \times SU(2)_R \rightarrow SU(2)_D$, $y = \pi R$ brane: $SU(2)_R \times U(1)_{B-L} \rightarrow U(1)_Y$. Without going into the details, for which we refer the readers to [32, 33], we mention that the W and Z boson masses, and the (S, T) parameters can be nicely fit in a warped scenario.

We highlight here two features which deserve attention.

(i) **Tension between unitarity and EWPT:** Recall that without a Higgs, unitarity violation would set in the SM at around a TeV. What is expected in the Higgsless scenario? Here, the exchange of KK states would retard the energy growth of the W_L - W_L scattering amplitude, *postponing* the violation of unitarity in a *calculable* way beyond a TeV. More specifically, $\Lambda \sim 3\pi^4 M_W^2 / (g^2 M_W^{(1)}) \sim 4$ TeV for $M_W^{(1)} \sim 1$ TeV. If we want to postpone the onset of unitarity violation even further, we have to decrease the $W^{(1)}$ mass. But this, in turn, increases the T parameter, implying a tension between unitarity and EWPT [34].

(ii) **LHC signature:** We deal with a specific signature here [35]. Consider the scattering channel $WZ \xrightarrow{W^{(1)}} WZ$. If $M_W^{(1)} \approx 700$ GeV, it turns out that $g_{WZV^1} \leq \frac{g_{WWZ} M_Z^2}{\sqrt{3} M_W^{(1)} M_W} \sim 0.04$. We then expect to see sharp resonance due to s channel mediation, with a striking feature of narrow width – see Fig. 7b.

VI Comparing Little Higgs/Composite Higgs/Higgsless scenarios

(i) **Little Higgs vs Composite:** Little Higgs models were introduced to solve the little hierarchy problem, but these models, as we saw before, are still quite fine-tuned. Moreover, the Higgs mass has a quadratic cutoff sensitivity at two-loop level. The composite Higgs does much better in both these aspects. It can have a proper UV completion all the way to M_{Pl} and the Higgs mass is finite at all order due to non-locality.

At an observational level, the composite Higgs model contains KK gluon (as it is dual to a 5d gauge theory), while a conventional little Higgs model does not have a KK gluon.

Fine-tuning in a composite model boils down to ensuring that $\xi = v^2/f^2 \sim 0.2$. The new resonances in composite models weigh around $g_\rho f$, where g_ρ could be as large as 4π . So the new resonances are heavy enough and their effects on EWPT normally die out. But there is a subtle point here which differentiates composite Higgs from the conventional little Higgs. The composite Higgs couplings to the gauge bosons are different from the corresponding SM couplings. These couplings pick up a factor ξ after the Higgs kinetic term is canonically normalized. As a result, the smooth cancellation of log divergence between the Higgs and gauge boson contributions to the S and T parameters does not hold any more, yielding an IR contribution $\sim \xi \ln(m_\rho/m_h)$ to the EWPT parameters [36]. This contribution is numerically very important, but a value of ξ fine-tuned to approximately 0.2 keeps the EWPT constraints under control. In conventional little Higgs models, the new heavy states weigh around gf , where g is the SM gauge coupling, and consequently the new resonances are not heavy enough. Therefore the little Higgs resonances pose a threat to overshoot the EWPT constraints.

(ii) **Composite vs Higgsless/Technicolor:** In the Technicolor model, QCD-like strong dynamics breaks electroweak symmetry *directly*. The 5d Higgsless model can be seen as dual to the ‘Walking Technicolor’. In a composite Higgs model, the strong sector does *not directly* break electroweak symmetry, but just delivers a composite pseudo-Goldstone boson, the Higgs, which gets a potential at one-loop and triggers electroweak breaking. This two-stage breaking, creating the parameter ξ , turns out to be a boon while facing EWPT constraints. The S parameter in a composite Higgs model is under control and, in fact, suppressed by the factor $\xi \sim 0.2$ compared to the value of S in the Higgsless model.

(iii) **Interpolation:** The composite Higgs scenario interpolates between the SM and the Higgsless (or, Technicolor) models in the two extreme limits of ξ [4]:

$$\text{SM} \xleftarrow{0 \leftarrow \xi} \text{Composite} \xrightarrow{\xi \rightarrow 1} \text{Higgsless/TC}$$

VII Conclusions and Outlook

1. All the BSM models we have considered are based on *calculability*. M_Z can be expressed in terms of some high scale parameters a_i , i.e. $M_Z = \Lambda_{\text{NP}} f(a_i)$, where $f(a_i)$ are calculable functions of physical parameters. The new physics scales originate from different dynamics in different cases: $\Lambda_{\text{SUSY}} \sim M_S$ (the supersymmetry breaking scale); $\Lambda_{\text{LH}} \sim f \sim F$ (the vev of $G \rightarrow H$ breaking); $\Lambda_{\text{Extra-D}} \sim R^{-1}$ (the inverse radius of compactification).

2. In supersymmetry the cancellation of quadratic divergence takes place between a particle loop and a sparticle loop. Since a particle and a sparticle differ in spin, the S, T, U parameters and the $Zb\bar{b}$ vertex correction can be kept under control, since new physics appear through loops. In the little Higgs scenario the cancellation occurs between loops with the same spin states. Such states can mix among themselves, leading to dangerous tree level contributions to the oblique parameters. This is the reason why a decoupling theory like supersymmetry is comfortable with EWPT, while a Technicolor-like non-decoupling theory faces a stiff confrontation.

3. We have a three-fold goal while building BSM physics: (i) unitarize the theory, (ii) successfully confront EWPT, and (iii) maintain as much naturalness as possible. The tension arises as ‘naturalness’ demands the spectrum to be compressed, while ‘EWPT compatibility’ pushes the new states away from the SM states.

4. Supersymmetric theories are getting increasingly fine-tuned with non-observation at LHC (having said that, we must realize that LHC direct searches do not apply for third generation matter superfields, and so a spectrum with inverted hierarchy is relatively less tuned). Naturalness in MSSM improved when we added extra singlets and additional terms in the superpotential. Although supersymmetry solves the big hierarchy problem by stabilizing the weak scale over many decades in energy scale, the little hierarchy problem continues to haunt and instigate the model builders to take bold, sometimes outrageous, steps for reducing the fine-tuning.

5. A light Higgs need not necessarily be elementary. It can very well be a composite object. Also, a narrow width of Higgs does not necessarily attest its elementarity. A light composite Higgs can very well have a narrow width. Just finding the Higgs would not settle this issue. We need to measure the Higgs couplings very precisely to know whether it is elementary or composite.

6. LHC is a ‘win-win’ discovery machine. If we find the Higgs, it will be a great discovery. But if we see only the Higgs and nothing else, we will definitely be disappointed as many of our questions would be left unanswered [37]. If, on the other hand, LHC confirms that there is no Higgs, it will be no less a discovery [38]. If the Higgs is not there, the new resonances which would restore unitarity in gauge boson scattering should show up, with a prior hint of excess events in $V_L V_L \rightarrow V_L V_L$ scattering. In that case we absolutely need the super-LHC, and eventually the ILC, to confidently establish the nature of the new resonances.

7. The excluded region in BSM parameter space is growing fast as LHC accumulate more and more data [39]. By the time we meet in the next Lepton Photon Conference in 2013, many of the possibilities discussed here may not perhaps be heard again! But who can rule out the possibility that completely new theoretical ideas inspired by yet unseen unexpected observations during the next two years would form the cornerstone for building the physics of the TeV scale?

Note added: On 13 December 2011 at CERN, the ATLAS and CMS experiments presented their update on the Higgs searches. The SM Higgs mass is now allowed in a narrower window: $115 < m_h^{\text{SM}} < 127$ GeV at 95% C.L. with a mild excess around 125 GeV. The main ethos of this talk remains unaffected by this observation. The only observation we can make at this stage is that if the Higgs mass is later confirmed to be around 125 GeV, some of the supersymmetric models we have discussed would require more tuning than before.

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References

- [1] G. Bhattacharyya, Rept. Prog. Phys. **74** (2011) 026201 [arXiv:0910.5095 [hep-ph]], and references therein.
- [2] G. F. Giudice, arXiv:0801.2562 [hep-ph].
- [3] R. Contino, arXiv:1005.4269 [hep-ph].
- [4] C. Grojean, PoSEPS -**HEP2009** (2009) 008 [arXiv:0910.4976 [hep-ph]].
- [5] See the text books on supersymmetry,
R.N. Mohapatra, “Unification and Supersymmetry: The Frontiers of quark-lepton physics,” Springer-Verlag, NY 1992; M. Drees,

- R. Godbole and P. Roy, “Theory and phenomenology of sparticles: An account of four-dimensional $N=1$ supersymmetry in high energy physics,” World Scientific (2004); H. Baer and X. Tata, “Weak scale supersymmetry: From superfields to scattering events,” Cambridge, UK: Univ. Pr. (2006). See also, R.K. Kaul, “Supersymmetry and supergravity,” in *Gravitation, gauge theories and the early universe*, 487-522, Ed. B.R. Iyer, Kluwer Academic Publishers (1989).
- [6] R. Barbieri and G. F. Giudice, Nucl. Phys. B **306** (1988) 63.
 - [7] A. Strumia, JHEP **1104** (2011) 073 [arXiv:1101.2195 [hep-ph]].
 - [8] U. Ellwanger, G. Espitalier-Noel and C. Hugonie, arXiv:1107.2472 [hep-ph].
 - [9] S. Cassel, D. M. Ghilencea, S. Kraml, A. Lessa and G. G. Ross, JHEP **1105** (2011) 120 [arXiv:1101.4664 [hep-ph]].
 - [10] For reviews on NMSSM, see U. Ellwanger, C. Hugonie and A. M. Teixeira, Phys. Rept. **496** (2010) 1 [arXiv:0910.1785 [hep-ph]]; M. Maniatis, Int. J. Mod. Phys. A **25** (2010) 3505 [arXiv:0906.0777 [hep-ph]].
 - [11] M. Drees, Int. J. Mod. Phys. A **4** (1989) 3635.
 - [12] G. G. Ross and K. Schmidt-Hoberg, arXiv:1108.1284 [hep-ph].
 - [13] S. A. Abel, S. Sarkar and P. L. White, Nucl. Phys. B **454** (1995) 663 [hep-ph/9506359].
 - [14] P. Lodone, JHEP **1005** (2010) 068 [arXiv:1004.1271 [hep-ph]].
 - [15] G. Bhattacharyya and A. Romanino, Phys. Rev. D **55** (1997) 7015 [hep-ph/9611243].
 - [16] M. Schmaltz and D. Tucker-Smith, Ann. Rev. Nucl. Part. Sci. **55** (2005) 229 [arXiv:hep-ph/0502182]; M. Perelstein, Prog. Part. Nucl. Phys. **58** (2007) 247 [arXiv:hep-ph/0512128]; M. C. Chen, Mod. Phys. Lett. A **21** (2006) 621 [arXiv:hep-ph/0601126].
 - [17] H. C. Cheng, arXiv:0710.3407 [hep-ph].
 - [18] R. K. Kaul, arXiv:0803.0381 [hep-ph].
 - [19] C. Csaki, J. Hubisz, G. D. Kribs, P. Meade and J. Terning, Phys. Rev. D **67** (2003) 115002 [arXiv:hep-ph/0211124]; J. L. Hewett, F. J. Petriello and T. G. Rizzo, JHEP **0310** (2003) 062 [arXiv:hep-ph/0211218]; M. C. Chen and S. Dawson, model,” Phys. Rev. D **70** (2004) 015003 [arXiv:hep-ph/0311032].
 - [20] I. Low, JHEP **0410** (2004) 067 [arXiv:hep-ph/0409025]; H. C. Cheng and I. Low, JHEP **0408** (2004) 061 [arXiv:hep-ph/0405243].
 - [21] J. Hubisz, P. Meade, A. Noble and M. Perelstein, JHEP **0601** (2006) 135 [arXiv:hep-ph/0506042].
 - [22] J. A. Casas, J. R. Espinosa and I. Hidalgo, JHEP **0503** (2005) 038 [hep-ph/0502066].
 - [23] T. Han, H. E. Logan, B. McElrath and L. T. Wang, Phys. Rev. D **67** (2003) 095004 [arXiv:hep-ph/0301040].
 - [24] G. Burdman, M. Perelstein and A. Pierce, Phys. Rev. Lett. **90** (2003) 241802 [Erratum-ibid. **92** (2004) 049903] [arXiv:hep-ph/0212228].
 - [25] M. Perelstein, M. E. Peskin and A. Pierce, Phys. Rev. D **69** (2004) 075002 [arXiv:hep-ph/0310039].
 - [26] K. Agashe, R. Contino and A. Pomarol, Nucl. Phys. B **719** (2005) 165 [hep-ph/0412089].
 - [27] S. Rychkov, arXiv:1109.1180 [hep-ph].
 - [28] M. Serone, New J. Phys. **12** (2010) 075013 [arXiv:0909.5619 [hep-ph]].
 - [29] R. Contino, Y. Nomura and A. Pomarol, Nucl. Phys. B **671** (2003) 148 [arXiv:hep-ph/0306259].
 - [30] J. R. Espinosa, C. Grojean and M. Muhlleitner, JHEP **1005** (2010) 065 [arXiv:1003.3251 [hep-ph]].
 - [31] R. Contino and G. Servant, JHEP **0806** (2008) 026 [arXiv:0801.1679 [hep-ph]].

- [32] For a review on Higgsless models, see C. Csaki, J. Hubisz and P. Meade, arXiv:hep-ph/0510275 (TASI lectures).
- [33] C. Csaki, C. Grojean, H. Murayama, L. Pilo and J. Terning, Phys. Rev. D **69** (2004) 055006 [arXiv:hep-ph/0305237].
- [34] R. Barbieri, A. Pomarol and R. Rattazzi, Phys. Lett. B **591** (2004) 141 [arXiv:hep-ph/0310285]; G. Cacciapaglia, C. Csaki, C. Grojean and J. Terning, Phys. Rev. D **70** (2004) 075014 [arXiv:hep-ph/0401160]; H. Davoudiasl, J. L. Hewett, B. Lillie and T. G. Rizzo, Phys. Rev. D **70** (2004) 015006 [arXiv:hep-ph/0312193].
- [35] A. Birkedal, K. Matchev and M. Perelstein, Phys. Rev. Lett. **94** (2005) 191803 [arXiv:hep-ph/0412278].
- [36] R. Barbieri, B. Bellazzini, V. S. Rychkov and A. Varagnolo, Phys. Rev. D **76** (2007) 115008 [arXiv:0706.0432 [hep-ph]]; J. R. Espinosa, C. Grojean and M. Muhlleitner, JHEP **1005** (2010) 065 [arXiv:1003.3251 [hep-ph]].
- [37] M. E. Peskin, “Summary of Lepton Photon 2011,” arXiv:1110.3805 [hep-ph] (In these proceedings).
- [38] R. Heuer, Talk in Lepton-Photon 2011, titled ‘Whither Colliders After LHC ?’ (In these proceedings).
- [39] H. Bachacou, Talk in Lepton-Photon 2011, titled ‘BSM Results from LHC’ (In these proceedings).